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Contributions of Vector-Meson-Dominance to Charmed Meson Production in Inelastic Neutrino and Antineutrino Interactions

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ABSTRACT

We discuss the contribution to neutrino-hadron interactions due to F^* dominance of the weak vector current, and we give quantitative estimates of the x- and y-distributions for the deep inelastic cross section at various energies. Other contributions, such as the nonresonant continuum and axial vector terms, have been qualitatively discussed as well. We compare our results with experimental data on single muon and dimuon production. We also calculate the diffractive, "elastic" F^* production cross section.

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Recently several authors[1-3] have discussed the diffractive production of a vector boson carrying a new quantum number such as "charm", in neutrino and antineutrino scattering. These considerations are motivated by the observation of anomalies in the invariant-mass distribution and the y-distribution (for small x) in inclusive antineutrino reactions,[4] and more recently by the observation of dimuon events at Fermilab.[5] They are based on the theoretical possibility that a new vector boson (F^*) might exist which couples to the charged weak-interaction current j_μ

$$\langle F^{*+}(\epsilon) \mid j_\mu(0) \mid 0\rangle = \frac{\mu^2}{\gamma} \epsilon_\mu$$

in the same way as the ρ^0, ω, ϕ and ψ do to the electromagnetic current, where μ is the mass of the vector boson, ϵ its polarization vector. We have pursued this possibility following the suggestion of Gaillard, et al.[6] We find that our conclusions and views are not entirely in accord with those of the aforementioned preprints in circulation.

We consider, therefore, the mechanism depicted in Fig. 1, in which the interaction of the W-boson with the nucleon is assumed to be vector-dominated by the F^*. (There may be an analogous contribution to the axial vector piece which we ignore for now.)

The total inelastic cross section due to the contribution from F^* dominance of the W-boson can be written as

$$\frac{d^2\sigma}{dxdy} = \frac{G_F^2 ME}{\pi} \frac{\cos^2\theta_C}{\gamma^2} \left(\frac{\mu^2}{Q^2+\mu^2}\right)^2 \frac{Q^2(1-x)}{2\pi} \Big[y^2\sigma_\perp +$$

$$2(\sigma_s + \sigma_\perp)(1-y-\frac{M}{2E}xy)/(1+\frac{2xM}{yE})\Big] \qquad (1)$$

where, as usual, $ME = p\cdot k$, $M\nu = p\cdot q$

$$q^2 = -Q^2 \qquad x = Q^2/2M\nu$$

$$y = \nu/E. \qquad \text{(Note: } xy = Q^2/2ME\text{)}.$$

Parity conservation by the strong interactions implies that the cross section for right- and left-handed vector bosons are equal, $\sigma_R = \sigma_L \equiv \sigma_\perp$ (i.e., $W_3 = 0$). Thus, the contribution to neutrino and antineutrino cross sections is equal. The cross section for a longitudinal F^* is denoted by σ_s.

On the mass shell, we expect the energy dependence of the total cross section $\sigma_\perp$ to be approximately constant (or slowly increasing) asymptotically. Duality diagrams suggest that the F^*N total cross section would be exchange degenerate like the KN, NN, ϕ N, or ψN cross sections. Thus, its asymptotic behavior would be manifested shortly above threshold, and it will be a good approximation to neglect the leading energy dependent corrections. Without a reliable theory to guide us, it is inherently ambiguous how we extrapolate the cross sections in Q^2. Following custom, we will assume simply that $\sigma_\perp$ is independent of Q^2. Further controversy surrounds the longitudinal cross section σ_s, and we shall treat $R = \sigma_s/\sigma_\perp$ as a parameter, giving it values typically observed in the analogous process of ρ electroproduction.

We note that in some models of longitudinal vector dominance,[7] R is proportional to Q^2 ($R = \xi Q^2/\mu^2$), so that this contribution (Eq. 1) <u>does</u> scale for $Q^2 >> \mu^2$. In this case, we find, in the scaling limit,

$$\frac{\pi}{G_F^2 ME} \frac{d\sigma}{dxdy} = \frac{\cos^2\theta_C}{\gamma^2} \frac{\xi\mu^2\sigma_\perp}{\pi} (1-x)(1-y) , \tag{2a}$$

and so

$$\frac{\pi}{G_F^2 ME} \sigma = \frac{\cos^2\theta_C}{\gamma^2} \frac{\xi\mu^2\sigma_\perp}{4\pi} . \tag{2b}$$

To get some idea of the magnitude of this cross section, we need to assume some values for the parameters. Based on SU(4) mass formulas,[6] we take $\mu^2 = 5\ \mathrm{GeV}^2$. Using SU(4) invariance and the observed value for γ_ρ in e^-e^+ annihilation,[8] we choose $\frac{\gamma^2}{4\pi} = \frac{1}{2}\frac{\gamma_\rho^2}{4\pi} = 1.3$. The additive quark model suggests that $\sigma_\perp$ would be half the sum of the ψp and ϕp cross sections, so we choose $\sigma_\perp = 6$ mb. If we choose $\xi = 0.25$, as in electroproduction, then the right-hand side of Eq. 2b is 0.09. We expect the asymptotic behavior of the cross section on neutrons to be the same, so we would get twice this value for the sum of the proton and neutron cross sections. Later, we shall discuss the evaluation of Eq. (1) at finite energies, both with and without assuming scaling behavior for R. By way of comparison, however, we note that typical experimental data[9] give values for the $\frac{\pi}{G_F^2 ME}\sigma \approx 0.5$ for neutrino scattering and about 1/3 of that for antineutrino scattering. This would mean that, above charm threshold, the neutrino cross section may increase by about 20 percent; the antineutrino cross section, by about 60 percent. (There might be a similar contribution due to the axial vector contribution.) On the other hand, if R were constant or falling as Q^{-2}, the F^* contribution would be smaller and confined to the regime $Q^2 \lesssim \mu^2$. However, since $\mu^2 >> m_\rho^2$, this effect differs from the electroproduction case, inasmuch as it may remain quite significant out to $Q^2 \approx 10\ \mathrm{GeV}^2$ or so.

Turning away from the issue of extrapolation off the mass shell, we want to explore what happens in F^*N collisions. Since we assume $\sigma_\perp = 6$ mb, the elastic cross section is $\sigma_{el} \cong \sigma_\perp^2/16\pi b$. Taking $b \approx 4\ \mathrm{GeV}^{-2}$, a value of the order observed for ψ electroproduction,[10] we find $\sigma_{el} \approx 0.5$ mb, so that $\sigma_{el}/\sigma_\perp = 8$ percent, an even smaller ratio than for noncharmed hadrons. In pp collisions, single diffraction dissociation is between 2 and 4 mb (for each proton) and double diffraction dissociation considerably smaller, at least until well beyond the ISR energy range.[11] This allows us to estimate that, in F^*N collisions, diffraction dissociation of the nucleon would be of order $\left(\frac{\sigma_T^{F^*p}}{\sigma_T^{pp}}\right)^2$ times, say 4 mb, giving only 0.1 mb. Assuming F^* dissociation is comparably small and double dissociation is negligible, we conclude that <u>almost 90 percent of the F^*N total cross section is nondiffractive inelastic.</u> Whether this leads to the production in the laboratory of a fast F^* or of some other charmed particle requires further assumptions. However, unless this percentage changes drastically as we go off the mass shell, diffractive F^* production seems likely to contribute at most 5 percent of the antineutrino cross section and even less for neutrino scattering.[12] (We will return to a discussion of the elastic contribution later.)

Of course, if R is less than linear in Q^2, then this process vanishes asymptotically. Obviously then, at finite energies, this cross section will concentrate at small x.

We recall that in the Bjorken scaling limit, for spin - 1/2 constituents, the scaling cross sections for neutrinos and antineutrinos are of the form[13]

$$\frac{d^2\sigma^{\nu}}{dxdy} = \frac{G_F^2 ME}{\pi} x\left[F_L^{\nu}(x) + (1-y)^2 F_R^{\nu}(x)\right] \tag{2a}$$

$$\frac{d^2\sigma^{\bar{\nu}}}{dxdy} = \frac{G_F^2 ME}{\pi} x\left[F_L^{\bar{\nu}}(x)(1-y)^2 + F_R^{\bar{\nu}}(x)\right] , \tag{2b}$$

where $F_L(F_R)$ refers to scattering from quarks (antiquarks). Consequently, for comparison with the scaling limit, it will be useful to consider the dimensionless quantity $\frac{d^2N}{dxdy} = \frac{\pi}{G_F^2 ME} \frac{d^2\sigma}{dxdy}$.

In Fig. 2, we plot this quantity from Eq. (1) for E = 50 GeV. (The various parameters take the values given above, $\mu^2 = 5$ GeV^2, $\sigma_\perp = 6$ mb, $\gamma^2/4\pi = 1.3$, $R = Q^2/4\mu^2$.) There is little qualitative or quantitative change for other choices for R such as R = 0.4 or R = 0. Note that the cross section peaks at small y. There may be a threshold effect, which must be superposed on these curves. For example, if charm threshold were 4 GeV, then the above formula would cut off for $W^2 = 2MEy(1-x) \leq 15\ \mathrm{GeV}^2$. For $x \leq 0.1$ and $E = 50$ GeV, this would occur for $y \leq 0.17$. (Arrows indicate the position for each curve.)

In Fig. 3, we indicate the energy dependence of dN/dy for $x \leq 0.1$. Below about 30 GeV, the curves increase monotonically while, above this energy, the cross section peaks somewhere in the region $0 \leq y \leq 0.5$ and falls from there to $y = 1$. The magnitude of dN/dy at fixed y undergoes a rather complicated energy dependence, rising as the energy increases to some maximum value from which it falls to its asymptotic value. The actual contribution will, at lower energies, be further complicated because of kinematical threshold effects as explained above.

It is difficult to know how to compare this effect to the scaling behavior which would be present below charm threshold. For $x \leq 0.1$, the Gargamelle data receives most of its contributions from such small values of Q^2 that one cannot hope to be in the scaling region. To get a rough idea of the magnitude of the effect, we note that at Gargamelle,[9] for $x = 0.1$; the momentum fraction carried by quarks is about 1.0; by antiquarks, 0.2. As $x \to 0$, we expect them to approach equal constants whose precise value is anybody's guess, say, 0.6, a rather generous fraction. So, a not unreasonable estimate of the noncharmed contribution to $\frac{dN}{dy}$ for $x \leq 0.1$ might be

$$\frac{dN^{\nu}}{dy} = 0.08 + 0.04\,(1-y)^2$$

$$\frac{dN^{\bar{\nu}}}{dy} = 0.04 + 0.08\,(1-y)^2 .$$

In Figs. 4a and 4b, we plot these curves and also indicate the sum of this noncharmed contribution with the curves presented in Fig. 3 for various energies. For $E = 40$ GeV, the total cross section for charm production for $x \leq 0.1$ is 13 percent of the neutrino cross section and 18 percent of the antineutrino cross section. At $E = 80$ GeV, these fractions increase slightly to 16 percent and 21 percent, respectively. (They are largest for E about 350 GeV, where they constitute 18 percent and 24 percent, respectively.) Although these magnitudes are quite a respectable fraction of the total cross section, the shape of the antineutrino distribution is not at all flat. Only if this additional effect were sharply increasing as $y \to 1$ could we hope to counterbalance the $(1-y)^2$ term. Unless, for some reason, threshold effects persisted to very high energies, it is impossible for this sort of model to reproduce the observed constant y distribution for antineutrinos. This is apparently true regardless of the overall normalization of the charm production cross section and reflects only that this cross section peaks for $Q^2 \lesssim \mu^2$.

There may be many other contributions to the production of charmed hadrons even within the framework of vector dominance. A discussion similar to the one above could be given for the axial-vector current.

However, the apparent absence of noncharmed axial vector counterpart of the ρ meson causes us to hesitate before doubling the preceding estimate.[14] However, even without axial-vector mesons, one would expect non-resonant contributions to both the axial-vector and vector currents. We will comment further on the magnitude of the continuum contribution, but first let us discuss the question of vector-axial vector interference, independently of charm and of whether axial vector mesons exist.

Asymptotically, this contribution depends on whether the pomeron can cause parity change in forward "elastic" scattering (Fig. 5). Recall that, in general, the structure function W_3, which reflects the vector-axial vector interference term, does not contribute to the longitudinal cross section but contributes to the right-handed and left-handed cross sections equally but with opposite sign. Even if the pomeron were a pole of positive parity, no further consequences for W_3 follow from parity conservation by the strong interactions. It has sometimes been suggested that, in meson transitions, a pomeron pole would change spin and parity together, according to $P(-)^{\Delta J} = +1$, where $P = \pm$ is the product of intrinsic parities of the mesons whose spins differ by ΔJ.[15] There seems to be no experimental or theoretical basis for the rule, but it seems logically <u>possible</u> that, if 1^+ mesons existed, the vector-meson-dominated contribution to W_3 would vanish. It seems unlikely that it could be true for the continuum, for models, such as the Deck effect, are known to violate the rule.[16] If we could neglect the continuum contribution for small Q^2, say,

$Q^2 < 1\ \mathrm{GeV}^2$, then a good test of this rule would be whether or not $W_3 \to 0$ as $W^2 \to \infty$ in neutrino and/or antineutrino reactions. Until axial vector mesons are found, however, the issue cannot be tested here and, in any case, has little to do with charm.

Aside from its effect on W_3, one wonders how the continuum contributions would modify the conclusions drawn above from consideration of resonances alone. In the generalized vector dominance model for electroproduction,[7] a simple choice for this contribution is able to reproduce the scaling data, at least for small x. ($x \lesssim 0.2$). A similar approach here would substantially increase the charm cross section for large Q^2, but it is difficult to predict precisely how. It is probably more convenient to consider the asymptotic behavior in the framework of the quark-parton model.[17] Above charm threshold, the structure functions $F_{L,R}^{\nu,\bar{\nu}}$ (Eq. 2) are given by[18]

$$F_L^{\nu} = u + d + 2s \qquad F_R^{\nu} = \bar{u} + \bar{d} + 2\bar{c}$$
$$F_L^{\bar{\nu}} = u + d + 2c \qquad F_R^{\bar{\nu}} = \bar{u} + \bar{d} + 2\bar{s} \quad .$$

(These are for the sum of proton and neutron). Experimentally, it is observed that, for $x \leq 0.1$, the cross sections are flat in y for both neutrinos and antineutrinos. Clearly, the only way this could result would be if, for $x \leq 0.1$, the nucleon were predominantly made of strange quarks, an implausible and untenable hypothesis.

It is therefore impossible to think of the data as scaling but simply undergoing a step function discontinuity at charm threshold. The details of nonscaling, i.e., energy dependent, contributions are quite model dependent, but we will comment on the qualitative features. First, we expect a threshold W_0 in the missing mass W for the production of heavy charmed particles. If this were at 5 GeV, say, then $W^2 = 2\ MEy\ (1-x) \geq 24\ GeV^2$. Thus, at any given energy E, the large y, small x region is favored. As discussed earlier, the cross section from Eq. (1) peaks for $Q^2 \approx \mu^2$, and the continuum contribution would presumably have some effective threshold $m_0 > \mu$. So we must have $Q^2 = 2\ ME\ xy \geq m_0^2$ which, for fixed E, favors large x and large y. Thus, in any case, the effects of new production would show up first at large y. This observation is in qualitative agreement with the data[4] which, e.g., shows $\langle y \rangle$ for antineutrinos increasing from 0.25 at low energies (E < 30 GeV) to about 0.4 at higher energies. The x dependence, on the other hand, is considerably more complicated depending on an interplay between the threshold in the missing mass W and the new mass scale affecting the Q^2 dependence.

The antineutrino distributions in y are, for $x \leq 0.1$, much flatter[4] than would be expected on the basis of a valence quark picture. Antiquarks in the nucleon would add a constant term to the scaling distribution, so, as remarked earlier, this is difficult to understand. We certainly don't anticipate that the valence quark contributions just disappear as energy increases! It must be that the sizeable $(1-y)^2$ term is compensated by a term increasing with y due to the threshold W_0 and possibly also the mass scales μ and m_0.

Because of antiquarks appearing for small x, we would have expected a $(1-y)^2$ term in the neutrino distributions for $x \leq 0.1$. As in the antineutrino case, this must be offset by an increasing y distribution due to threshold effects. If there is a greater probability for finding quarks rather than antiquarks in the nucleon, even for $x < 0.1$, we would expect the coefficient of $(1-y)^2$ to be larger for antineutrinos than for neutrinos. Thus, in a model such as the one under consideration, making equal contributions to the neutrino and antineutrino cross sections, it would not be possible to compensate exactly in each case. Whether the introduction of a nonzero W_3 could change this conclusion, we have not investigated. We have, however, investigated models like those discussed in Ref. 4 and, with an appropriate choice of parameters, it is possible to accommodate the data at its present modest level of accuracy.

This completes our discussion of the total cross section; however, because of its distinctive experimental signature, we conclude with a discussion of the elastic F^*N contribution (Fig. 6). (As we pointed out earlier, the elastic F^*N cross section is apt to be a small fraction (< 10 percent) of the total F^*N cross section.) To this end, we may simply replace $\sigma_\perp$ in Eq. (1) by σ_{el}. However, there is one effect of extrapolating the initial F^* off the mass shell which is likely to be quite important, viz., the minimum momentum transfer allowed. So we multiply σ_{el} by $\exp(b\, t_{min})$, where, in the Bjorken limit,

$t_{min} \simeq \frac{-M^2x^2}{1-x} \left[1 + \frac{1}{2\,MEy}\left(M^2 + \frac{\mu^2}{x}\right)\right]$. Thus this exponential forces the cross section to be concentrated at small x. As indicated earlier, we take $b = 4\ \mathrm{GeV}^{-2}$ and $\sigma_{el} = 0.5$ mb. We take $R = (\sigma_s/\sigma_\perp)_{el} \approx Q^2/16\,\mu^2$. We will assume σ_{el} to be energy-independent. (Because the reggeons are exchange-degenerate, their contribution to the elastic amplitude is real. Hence the pomeron-reggeon interference term vanishes, and the leading correction to the asymptotic behavior of $\sigma_\perp$ is of order W^{-2}, which we neglect) In Fig. 7a, we plot the resultant distribution at 50 GeV as a function of y for various values of x. (The curves are terminated for small y, where W falls below kinematic threshold.) Notice that, because of the exponential factor, this decreases much more rapidly with x than does the total cross section (Fig. 2). Nevertheless, the cross section for $x \geq 0.1$ exceeds that of $x \leq 0.1$ at all energies. (It isn't clear to us whether the total cross section should not also show some t_{min} effect, but it would seem to be more model dependent and has been subsumed in our ignorance of the extrapolation off the mass shell.) In Fig. 7b, we plot dN/dy for several energies. This is, of course, a small fraction of the total cross section. For example, the area under these curves gives $\frac{\pi}{G_F^2\,ME}\,\sigma$ varying from $(2.6)(10^{-3})$ at 20 GeV to a maximum of $(4.9)(10^{-3})$ around 200 GeV. By comparison, the observed total cross section[4] appears to increase linearly with energy with $\pi\,\sigma_{TOT}/G_F^2\,ME \approx 0.41$ for neutrinos and one-third of that for antineutrinos. Thus, these elastic events make up at most about 1 percent of the total observed rate at any energy. However, their distinctive signature makes them detectible in a facility such as the 15 foot bubble chamber with the external muon identifier at Fermilab. Triggering on the muon and requiring that the only slow particle in the chamber be the recoil nucleon or nucleus, one could pick out the diffractive events.

To summarize, we have estimated the contribution to neutrino interactions due to F^* dominance of W-boson exchange. The magnitude is on the order of 20 percent (60 percent) of the neutrino (antineutrino) cross section. With a branching to leptons of only 5 to 10 percent, this could easily account for the rate observed for oppositely charged dimuons.[5] In the conventional charm scheme, there would be no dimuon events of the same sign directly, although some may result indirectly from mixing between neutral mesons of opposite charm (D^0-$\bar{D}^0$ mixing). However, it is difficult to believe the signal would be as large a fraction of the opposite sign events as appears to have been observed.[5] Moreover, we have seen that it is difficult, by the mechanism alone, to accommodate the "anomaly" observed in the y-distribution for antineutrinos. When added to the conventional noncharmed background, the effect does tend to flatten out the y distribution. However, because it adds the same to the neutrino cross section, it is difficult to see how one could obtain distributions nearly constant in y for both neutrinos and antineutrinos. It is hard to say whether other contributions, including vector-axial vector interference, could explain the data, although at its present level of accuracy, it seems possible.

Finally, we calculated "elastic" F^* production, where rate is small (about 1 percent of the total cross section) but whose signature is quite distinctive. Regardless of whether the charm hypothesis is correct, something new is happening and such events will provide information on the "anomalous" piece of the weak current.

Note Added: After complection of this manuscript, we received another preprint[19] on the same subject. Numerical differences from us arise from their inclusion of the A_1 and F_A^* as axial vector mesons (a factor of two) and their assumption that the transverse cross section is larger (by a factor of 1.5) than the total cross section for charm production. Although including axial mesons, they neglect W_3. Although we generally agree with their qualitative remarks, we think it more likely that, if anything, the continuum would account for the bulk of the anomaly than the adjustment of threshold and overall normalization required by their fits to the data.

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[12] Our estimates agree with those of Ref. 1 where, despite the title to the contrary, the authors calculate the total cross section and not just the diffractive contribution. The numerical estimates of the elastic contribution presented in Ref. 3 are too large, apparently because of an incorrect estimate of damping due to the minimum momentum transfer effect. (J. Pumplin, private communication.)

[13] See, e.g., S. L. Adler, review presented at NAL Topical Conference on Neutrino Physics, Batavia, March 1974 (NAL-Conf-74/39-THY, unpublished.)

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[16] G.C. Fox, private communication. We are grateful to Professor Fox for sharing his understanding of the controversy surrounding this subject.

[17] It is possible that a model in which the continuum contribution scales is simply an alternate (dual) description of the same physics.

[18] See, e.g., G. Altarelli, N. Cabibbo, L. Maiani, and R. Petronzio, Phys. Letters 48B, 435 (1974).

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FIGURE CAPTIONS

Fig. 1 F^* dominance of the weak vector current.

Fig. 2 Contribution from F^* dominance to $d^2N/dxdy$ at 50 GeV for (a) $x \geq 0.1$ and (b) $x \leq 0.1$. The position of a threshold at $W = 4$ GeV is indicated by the arrow above each curve.

Fig. 3 Energy dependence of $\frac{dN}{dy}$ $(x \leq 0.1)$.

Fig. 4 Total cross sections for (a) neutrino and (b) anti-neutrino scattering at $E = 40, 80, 200$ GeV.
- - - - - noncharmed model described in text.
———— noncharmed plus charmed cross section.

Fig. 5 Pomeron contribution to V-A interference.

Fig. 6 Elastic F^*N scattering.

Fig. 7 Contribution of elastic F^*N Scattering to (a) $d^2N/dxdy$ at 50 GeV and (b) dN/dy for several energies.

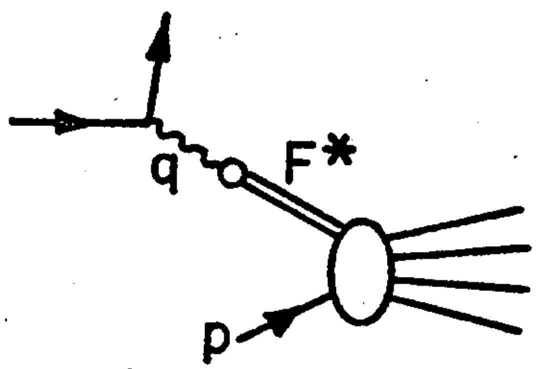
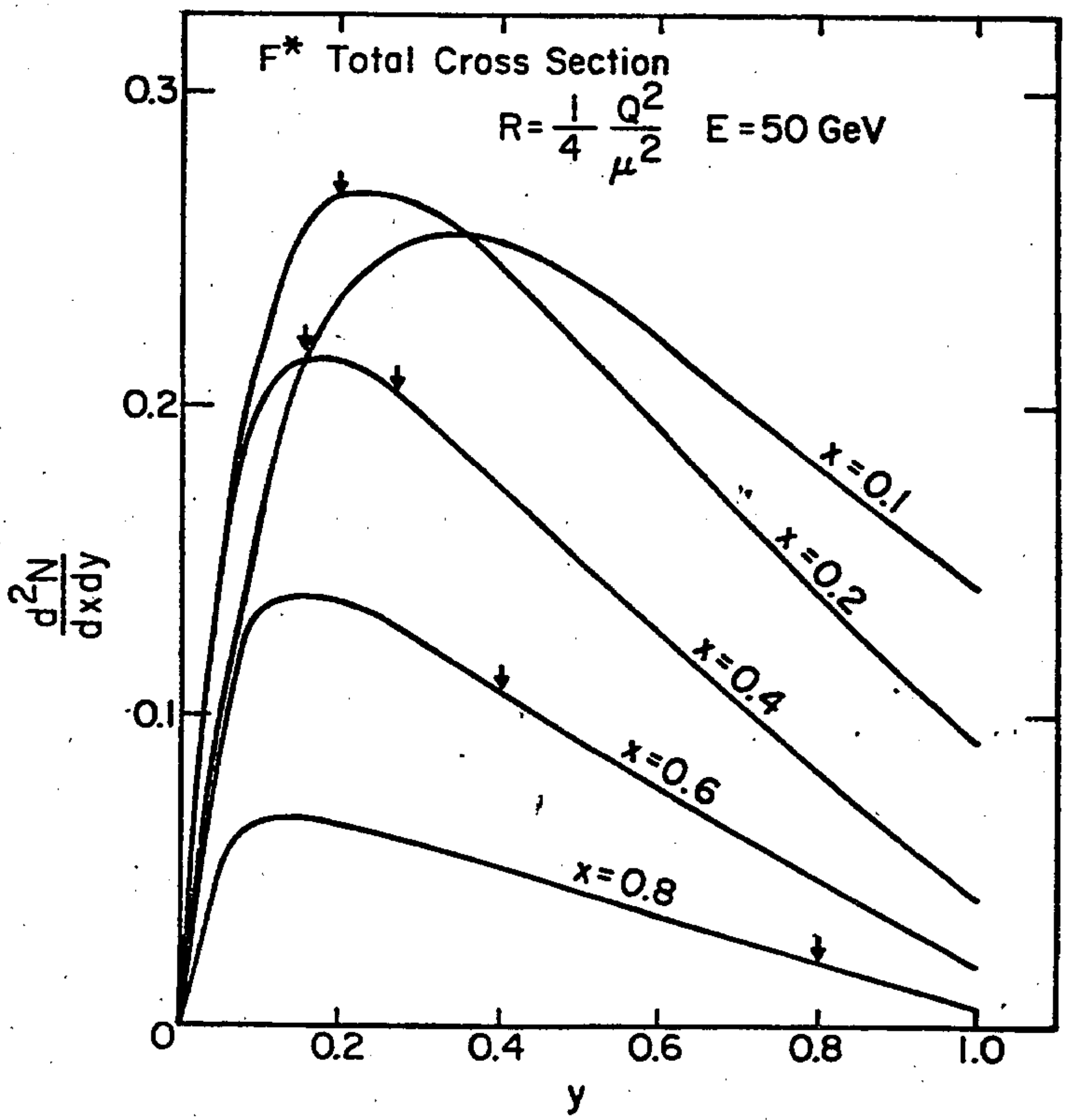
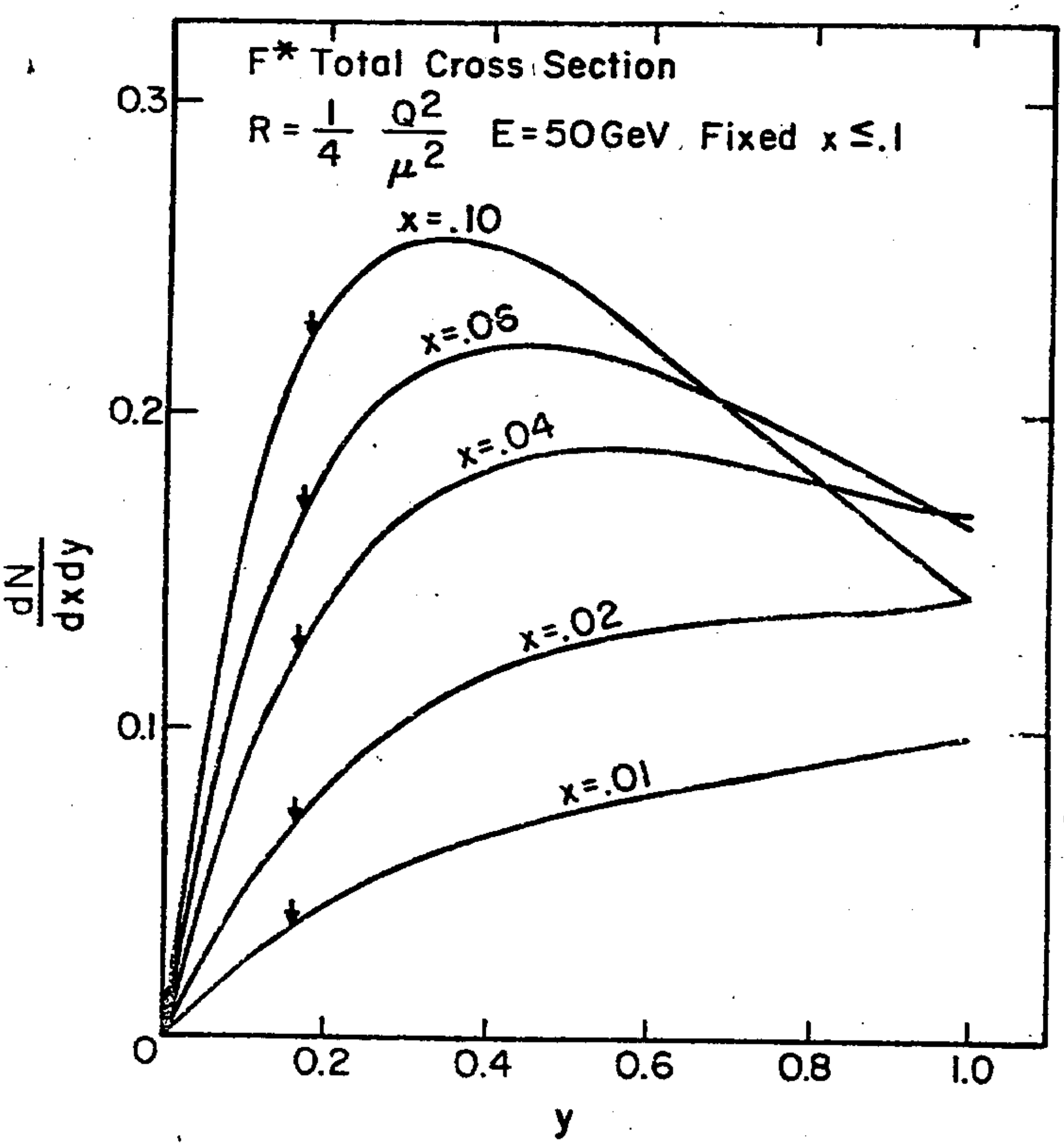


Fig. 1



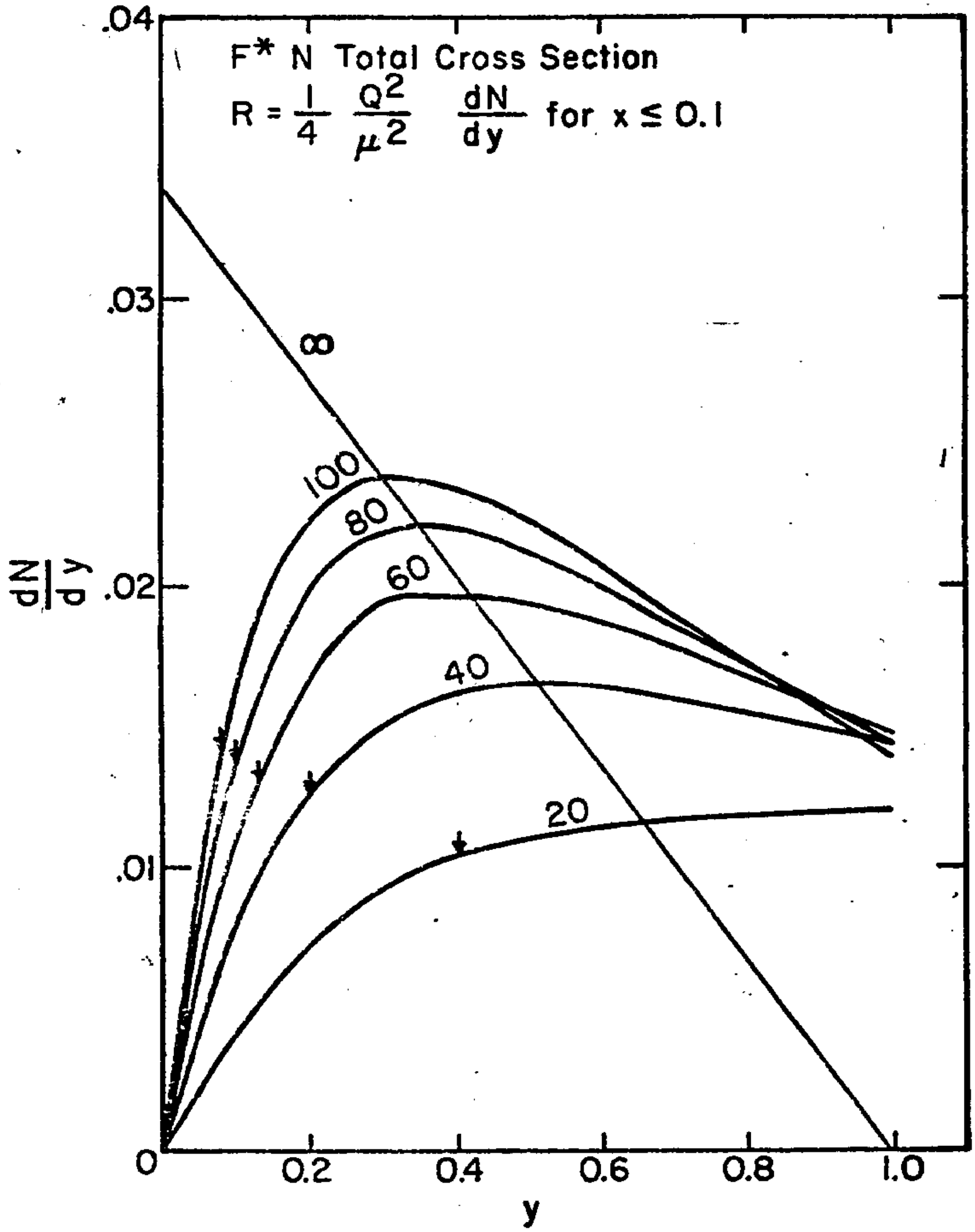
F* Total Cross Section
R = 1/4 Q²/μ²   E = 50 GeV
0.3
0.2
0.1
0
d²N/dxdy
x = 0.1
x = 0.2
x = 0.4
x = 0.6
x = 0.8
0.2
0.4
0.6
0.8
1.0
y

Fig. 2a



F* Total Cross Section
R = 1/4 Q2/μ2 E=50GeV, Fixed x ≤.1
x=.10
x=.06
x=.04
x=.02
x=.01
0.3
0.2
0.1
0
0.2
0.4
0.6
0.8
1.0
dN/dxdy
y

Fig. 2b



F* N Total Cross Section
R = 1/4 Q2/μ2 dN/dy for x ≤ 0.1
.04
.03
.02
.01
0
dN/dy
∞
100
80
60
40
20
0.2
0.4
0.6
0.8
1.0
y

Fig. 3

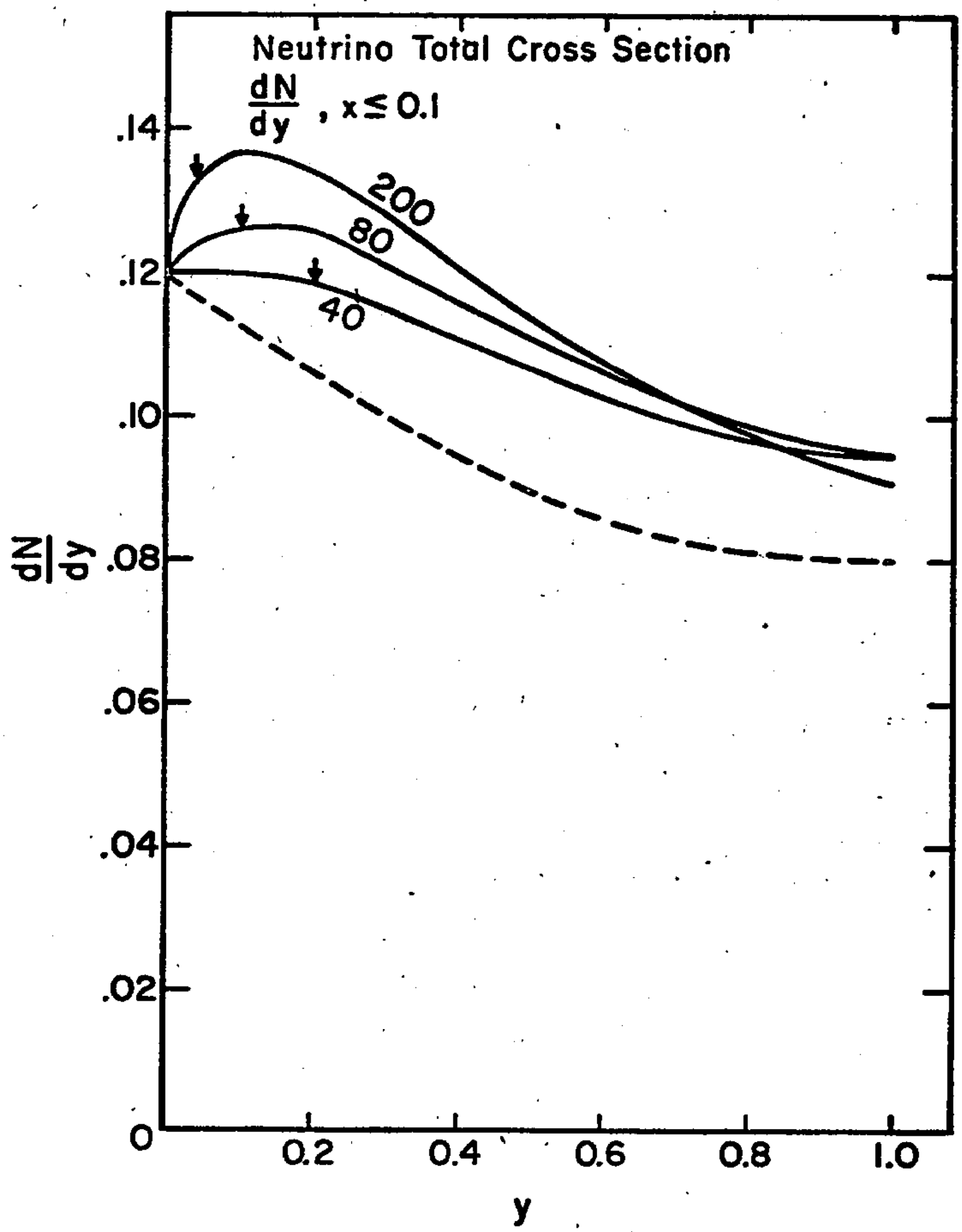
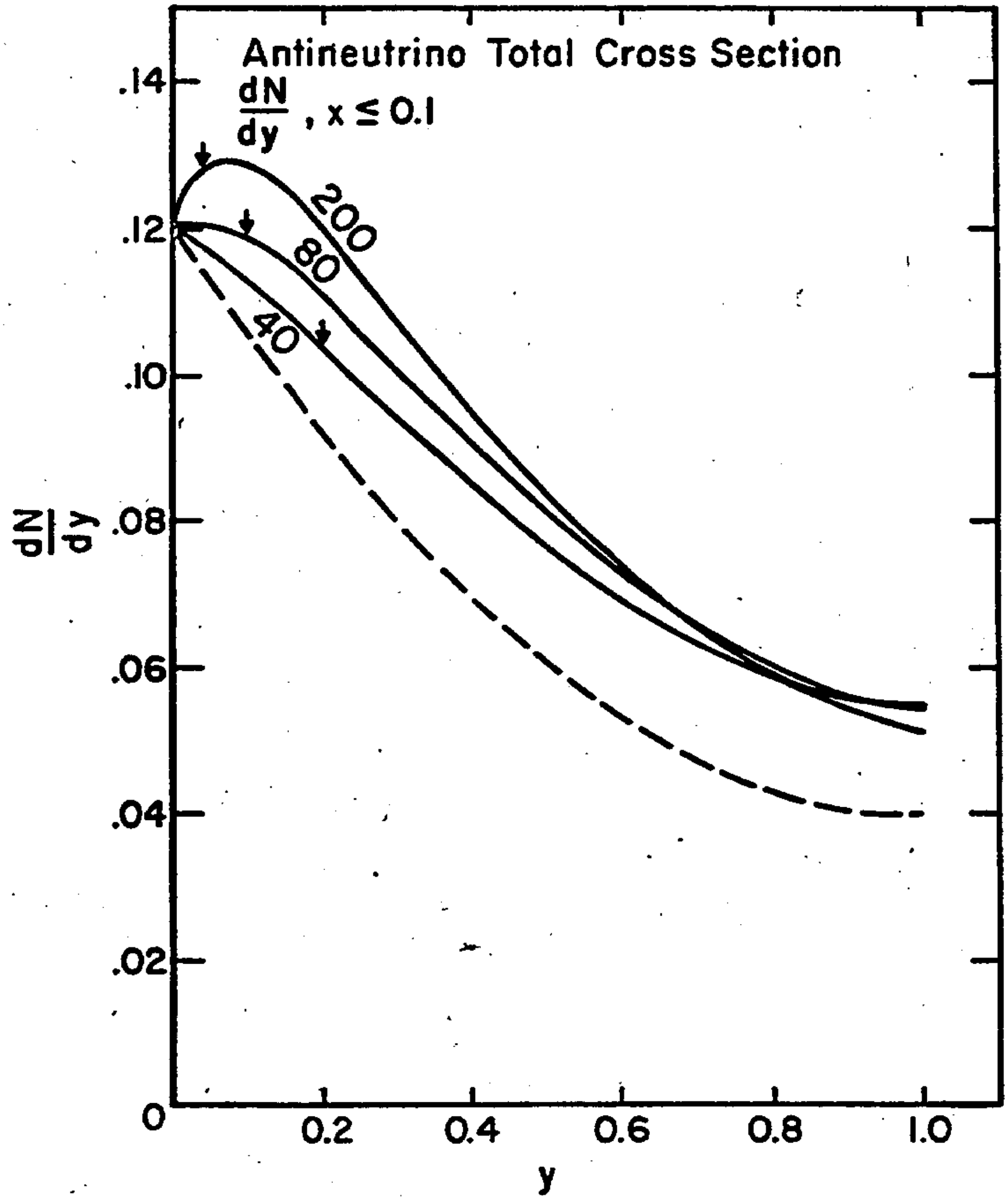

Neutrino Total Cross Section
dN/dy , x ≤ 0.1
200
80
40
.14
.12
.10
.08
.06
.04
.02
0
dN/dy
0.2
0.4
0.6
0.8
1.0
y

Fig. 4a



Antineutrino Total Cross Section
dN/dy , x ≤ 0.1
200
80
40
.14
.12
.10
.08
.06
.04
.02
0
dN/dy
0.2
0.4
0.6
0.8
1.0
y

Fig. 4b

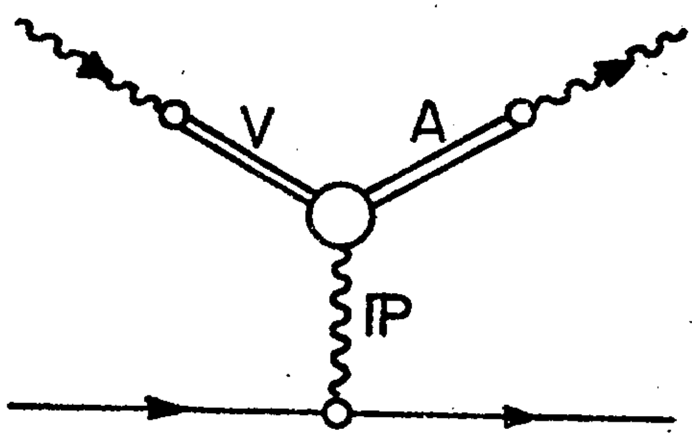


Fig. 5

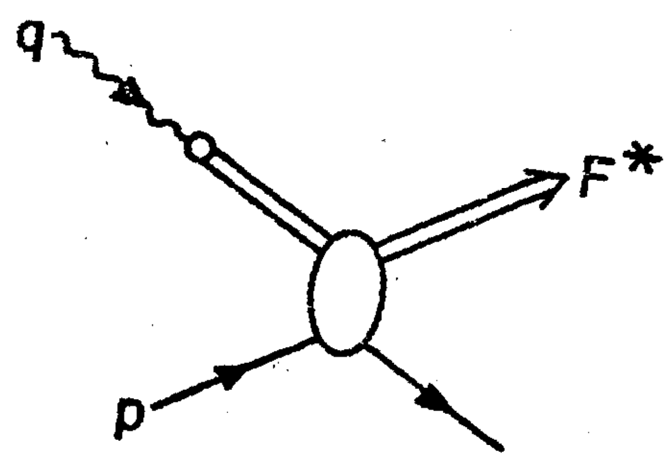


Fig. 6

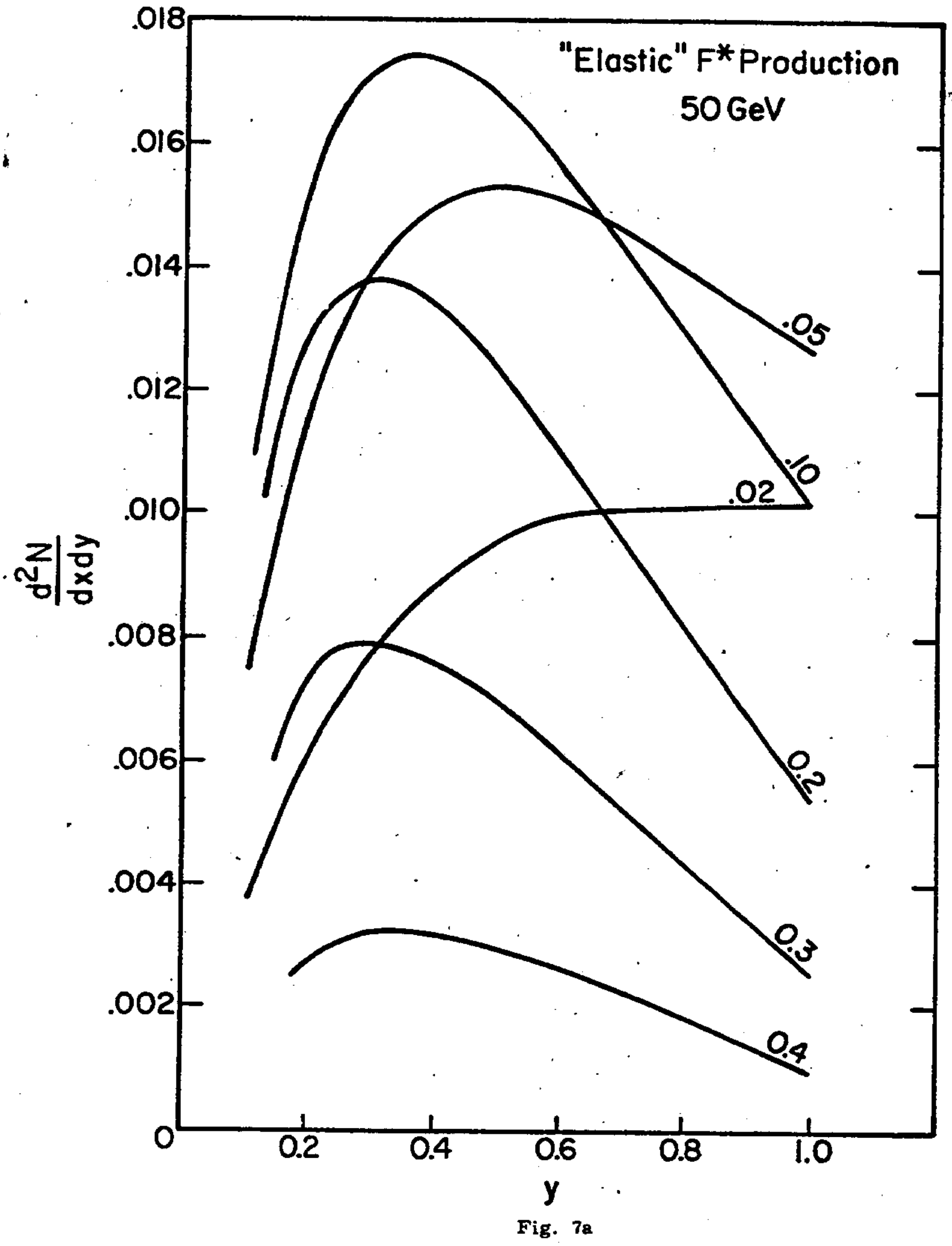
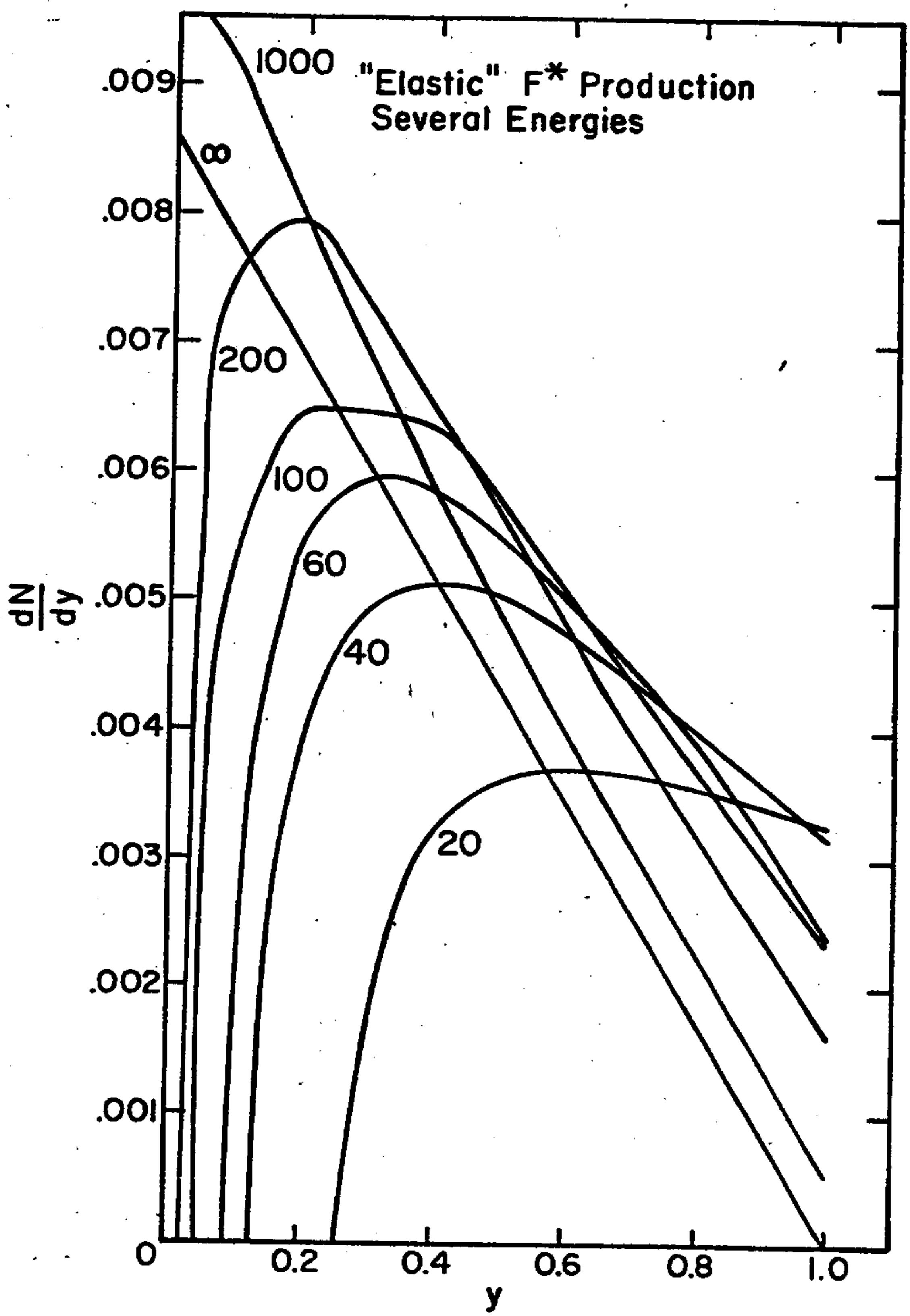

"Elastic" F* Production
50 GeV
.018
.016
.014
.012
.010
.008
.006
.004
.002
0
$\frac{d^2N}{dx\,dy}$
.05
.10
.02
0.2
0.3
0.4
0.2
0.4
0.6
0.8
1.0
y

Fig. 7a



"Elastic" F* Production
Several Energies
1000
∞
200
100
60
40
20
$\frac{dN}{dy}$
.009
.008
.007
.006
.005
.004
.003
.002
.001
0
0.2
0.4
0.6
0.8
1.0
y

Fig. 7b